

Phase Sensitivity of a Mach-Zehnder Interferometer

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The best performance of a Mach-Zehnder interferometer is achieved with the input state $|N_T/2 + 1\rangle|N_T/2 - 1\rangle + |N_T/2 - 1\rangle|N_T/2 + 1\rangle$, being N_T the total number of atoms/photons. This gives: i) a phase-shift error confidence $C_{68\%} = 2.67/N_T$ with ii) a *single* interferometric measurement. Different input quantum states can achieve the Heisenberg scaling $\sim 1/N_T$ but with higher prefactors and at the price of a statistical analysis of *two or more* independent measurements.

Introduction. The central goal of interferometry is to estimate phase shifts with the highest sensitivity given a finite resource of atoms/photons. There are several possible interferometric configurations, which are generally reducible to the “drosophila” Mach-Zehnder (MZ), see Fig.(1). The MZ transforms a two-mode input state as $|\psi\rangle_{out} = e^{-i\hat{J}_y\theta}|\psi\rangle_{inp}$ [1], where the generator of the phase translation θ is the y angular momentum component of the three-dimensional rotation group $\hat{J}_x = (\hat{a}^\dagger\hat{b} + \hat{a}\hat{b}^\dagger)/2$, $\hat{J}_y = (\hat{a}^\dagger\hat{b} - \hat{a}\hat{b}^\dagger)/2i$, $\hat{J}_z = (\hat{a}^\dagger\hat{a} - \hat{b}^\dagger\hat{b})/2$ (\hat{a}, \hat{b} are bosonic annihilation operators for the two input quantum modes). How precisely can the unknown phase shift θ be estimated from the output $|\psi\rangle_{out}$? It is well known that, if one of the two input arms of the interferometer is fed with vacuum, the phase measurement uncertainty is bounded by the standard quantum limit (or “shot noise”) $\Delta\theta \geq 1/\sqrt{N_T}$, where N_T is the total/average number of atoms in the input state. The first breakthrough came in the early ’80, when Caves showed that it is possible to reach sub shot-noise sensitivities after squeezing the vacuum of the unused port of the interferometer [2]. Sub shot-noise was experimentally demonstrated few years later [3]. This inaugurated a large body of literature proposing new states and optimal performances [1, 4, 5, 6, 7, 8] to push the sensitivity up to the Heisenberg limit $\sim 1/N_T$ [9]. Paradoxically, however, fewer efforts have been devoted to the correct inference of phase measurement uncertainties. This fact went (and it is still going) almost unnoticed, despite of the seminal works of Helstrom and Holevo [10] and Lane, Braunstein and Caves [11]. Apart from a few noticeable exceptions, in the current literature the phase measurement uncertainty is still retrieved from a simplified error propagation theory, $\Delta\theta = (\Delta\hat{A})/|\frac{\partial\langle\hat{A}\rangle}{\partial\theta}|$, where $\langle\hat{A}\rangle$ is the average of a generic phase dependent observable $\hat{A}(\theta)$ and $\Delta\hat{A} = \langle\hat{A}^2\rangle - \langle\hat{A}\rangle^2$ is the mean square fluctuation. Unfortunately, this simple approach can give correct estimates only when the probability distributions are Gaussian [12], which is seldom the case in quantum interferometry. Moreover, optimized phase estimations often require the distribution of the available finite resource of N_T particles among p independent interferometric experiments, each with $N = N_T/p$ particles. This optimization cannot be analyzed with the linear error propagation formula, which only considers the dependence on the num-

ber of particles in each experiment N and not on the total number of particles N_T . As a matter of fact, even if the Mach-Zehnder has been extensively studied, there is not published literature about the exact dependence of phase uncertainties on the total (or mean) number of particles of any input state.

Matter waves Mach-Zehnder interferometers have been realized with cold atoms [13] and electrons [14]. Evidences of non-classical light properties after a beam splitter have been experimentally showed in [15]. Entangled “NOON” states have been created with few photons [16] and ions [17] with applications in quantum lithography [18]. Current efforts are devoted to the creation of the MZ with dilute Bose-Einstein condensates (BEC). BEC offers almost unique possibilities to create the large intensity, highly squeezed waves needed for Heisenberg interferometry. Recently, double-slit [19] as well as Michelson interferometers [20] have been experimentally demonstrated. By trapping BEC in two wells and in periodic potentials, the beam splitters and mirrors are replaced by the dynamical tailoring of interwell barriers and magnetic wells. Fock and number squeezed condensates have been realized in [21].

In the following, we study a MZ interferometer fed by

$$|\psi\rangle_{inp} = \frac{1}{\sqrt{2}}(|j+m\rangle_a|j-m\rangle_b + |j-m\rangle_a|j+m\rangle_b), \quad (1)$$

where $N = 2j$ and $2m$ are the total and the relative number of particles at the a and b input ports, respectively, cfr. Fig. (1). We rigorously calculate error confidences within a Bayes framework [10]. We first consider the twin-Fock ($m = 0$) introduced by Holland and Burnett [5]. We then study a “NOON” state ($m = j$), showing that it cannot overcome shot-noise sensitivity. We will finally demonstrate that the best Mach-Zehnder performance is obtained when $m = 1$. This state can be created by constructing a twin-Fock state followed by the measurement of one particle [22].

Bayes analysis. Quantum Mechanics provides the conditional probability $P(\mu|j, \theta)$ to measure, at the output ports, the relative number of atoms $\mu = (N_1 - N_2)/2$, given $2j = N_1 + N_2$ particles and the unknown phase shift θ . With the general state Eq.(1) we have

$$P(\mu|j, \theta) = |\langle j, \mu | \psi \rangle_{out}|^2 = \frac{1}{2} |d_{\mu, m}^j(\theta) + d_{\mu, -m}^j(\theta)|^2, \quad (2)$$

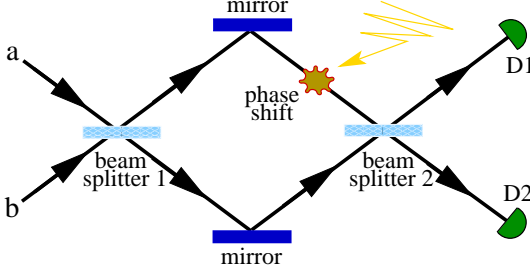


FIG. 1: Schematic representation of the Mach-Zehnder interferometer. Atoms/photons enter the a and b input ports, mix and recombine in the beam splitters and are finally detected in D1 and D2. The phase shift is inferred from the number of atoms/photons measured in each output port.

where $|j, \mu\rangle \equiv |j + \mu\rangle_{D1} |j - \mu\rangle_{D2}$, being D1 and D2 the detectors at the output ports (see Fig. (1)), and $d_{\mu, \nu}^j(\theta)$ being the angular momentum matrices [23]. The goal is to estimate θ after detecting a certain value μ . According to the Bayes theorem, $P(\theta|j, \mu)P(\mu) = P(\mu|j, \theta)P(\theta)$, where $P(\theta)$ and $P(\mu)$ take into account any further *a priori* information about the real value of the phase shift and the output measurement μ . We assume complete ignorance, so $P(\theta) = H(\pi/2 - |\theta|)$, being $H(x)$ the Heaviside step function, while $P(\mu)$ is fixed by the normalization. We notice that, with the state Eq.(1) as input of the MZ, it is possible to estimate only the absolute value of phase shifts in the interval $[0, \pi/2]$. In most cases, the best interferometric performance requires the statistical analysis of several independent measurements. Once obtained the results $\mu_1 \dots \mu_p$, the phase probability distribution reads $P_p(\phi|j, \mu_1 \dots \mu_p) \propto \prod_{k=1}^p P(\phi|j, \mu_k)$. Averaging over all the possible p -uple $\mu_1 \dots \mu_p$ we find the phase distribution for fixed θ and j :

$$P_p(\phi|j, \theta) = \sum_{\{\mu_1 \dots \mu_p\}} P_p(\phi|j, \mu_1 \dots \mu_p) P_p(\mu_1 \dots \mu_p|j, \theta), \quad (3)$$

where $P_p(\mu_1 \dots \mu_p|j, \theta) = \prod_{k=1}^p P(\mu_k|j, \theta)$ is provided by quantum mechanics. As phase estimator, we choose the value of the phase $\bar{\phi}_p(j, \theta)$ corresponding to the maximum of the probability distribution Eq.(3). The phase uncertainty is estimated as confidence, namely the γ -probability that the real value of the phase shift is within a given interval c_γ around in $\bar{\phi}$: $\gamma = \int_{\bar{\phi}-c_\gamma}^{\bar{\phi}+c_\gamma} d\phi P_p(\phi|j, \theta)$.

In the following we consider the case of a null phase shift $\theta = 0$. The probability distribution Eq.(2) reduces to $P(\mu|j, 0) = |\delta_{\mu, m} + \delta_{\mu, -m}|^2/2$, where $\delta_{\alpha, \beta}$ is the Kronecker delta. Therefore, we have two possible results for the measurement of the relative number of particles: $\mu = +m$, and $\mu = -m$. Because of the symmetry properties of the MZ interferometer and of the state Eq.(1), we have that $P(\phi|j, m) = P(\phi|j, -m)$, and the sum Eq.(3) reduces to $P_p(\phi|j, 0) \approx |d_{m, m}^j(\phi) + d_{m, -m}^j(\phi)|^{2p}$. When $\theta \neq 0$ the probability of obtaining values of $\mu \neq \pm m$ increases. In this case we might need several measurements in order to explore all the relevant probability distribu-

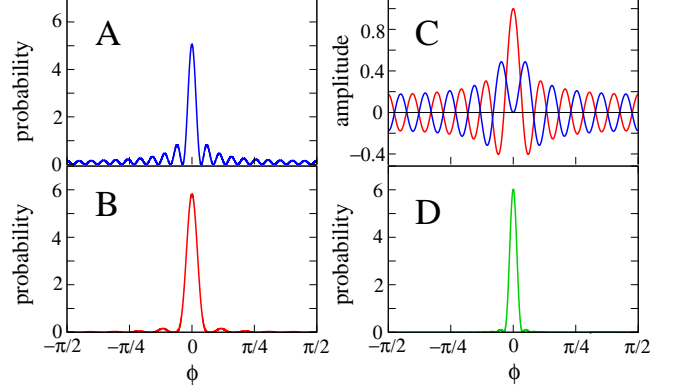


FIG. 2: A,B: Probability distribution of the $m = 0$ twin-Fock state with A) $N = 40$ particles and a $p = 1$ measurement, (A); $N = 20$ particles and $p = 2$ measurements, (B). By combining the independent measurements it is possible to strongly reduce the weight of the tail of the distribution with respect to the central peak. C,D: Amplitude $d_{1,1}^j(\phi)$ (red line), and $d_{1,-1}^j(\phi)$ (blue line), for $N = 40$ particles ($j = 20$) [23], (C). The tails of the amplitudes oscillate out of phase and interfere destructively, so to enhance the central peak of the probability distribution $P(\phi|j, 0) \sim |d_{1,1}^j(\phi) + d_{1,-1}^j(\phi)|^2$, (D).

tion, so the sensitivity of the interferometer for a given number of particles is, in general, reduced [24].

Mach-Zehnder interferometer with twin-Fock States. We first consider the input state $|\psi\rangle_{inp} = |j\rangle_a |j\rangle_b$ ($m = 0$ in Eq.(1)).

The probability of measuring a relative particle number μ at the output port, given a phase shift θ , is [25]: $P(\mu|j, \theta) = \frac{(j-\mu)!}{(j+\mu)!} [P_j^\mu(\cos \theta)]^2$, where $P_j^\mu(x)$ are the Associated Legendre Polynomials. When $\theta = 0$ the probability $P(\mu|j, \theta)$ reduces to a single peak centered in $\mu = 0$. Therefore, repeating the experiment p times we have $P_p(\phi|j, 0) \approx [P_j^0(\cos \phi)]^{2p}$. With a single measurement ($p = 1$), the 68%-confidence scales as $\sim 1/N_T^{1/3}$ while, as discussed in [26], the mean square phase fluctuation scales as $\sigma \sim 1/\ln N_T$. In both cases the phase uncertainty is worse than shot noise [27]. This problem is generally overlooked when the sensitivity is calculated from the simplified error propagation formula, or looking at the width of the central peak of the Legendre polynomials, which indeed shrink as $1/N_T$. The large tails of the distribution contain most of the probability and cannot be ignored, see Fig.(2,A). This clarifies the need of considering several independent measurements, which increases the relative weight of the central peak with respect to the tails, see Fig. (2,B). As discussed in [26], the mean square phase fluctuation is minimized combining $p = 4$ measurements.

In Fig. (3,A) we study different values of confidence for the twin-Fock state distribution as a function of the number p of independent measurements and a fixed total number of atoms N_T . We see that the bigger the confidence, the larger is the number of experiments, p , necessary to reach the minimum. It is important to no-

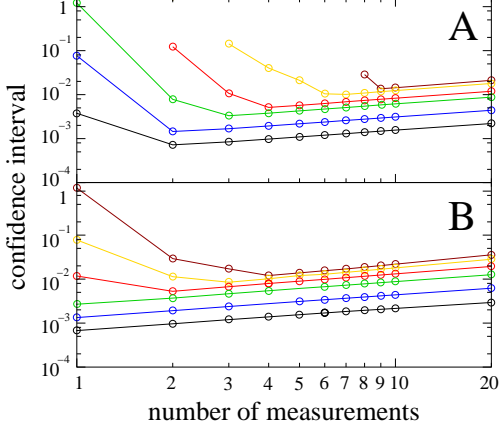


FIG. 3: Confidence interval as a function of the number of independent measurements p for A) the twin-Fock $m = 0$ and B) the twin $m = 1$ states. Here $N_T = 2000$ particles. In both plots, from the bottom line to the top one, the black line is the confidence at 38.29%, the blue at 68.27%, the green at 95.45%, the red at 99.73%, the yellow at 99.994%, and the brown at 99.99994%. These correspond to $\sigma/2, \sigma, 2\sigma, 3\sigma, 4\sigma$ and 5σ when the distribution is Gaussian.

tice that the number of measurements needed to reach the minimum for a given confidence does not depend on the total number of atoms N_T . The main results of our analysis can be summarized as follows: i) with a single measurement the MZ interferometer gives a sensitivity worse than the shot noise, ii) the highest sensitivity is reached with *two* independent measurements, for the 68% confidence, and iii) with *three* measurements for the 95% confidence:

$$\Delta\theta_{68\%} = \frac{2.915}{N_T}, \quad p = 2; \quad \Delta\theta_{95\%} = \frac{6.654}{N_T}, \quad p = 3. \quad (4)$$

The number of measurements needed to reach higher confidences is shown in Fig.(3,A). In Fig.(4,A) we summarize the scaling of the phase uncertainties reached with the twin-Fock state, as a function of N_T .

How does the sensitivity scale with the number of measurements, p , given a fixed number of atoms, N , for each experiment? The Cramer-Rao theorem does not allow a better scaling than $1/\sqrt{p}$, and we have verified numerically that the limit is reached when $p > 4$, obtaining the mean square fluctuation $\Delta\theta \approx 2\sqrt{p}/N_T$. The $1/p$ scaling claimed in [25] was probably an artifact of low statistics.

Mach-Zehnder interferometer with NOON state. We now consider the state (1) with $m = j$. If $\theta = 0$ the only possible outcomes are $\mu = \pm j$. According to the Bayes theorem, we obtain the probability distribution $P_p(\phi|j, 0) \approx [\cos^{2j}(\phi/2) + \sin^{2j}(\phi/2)]^{2p}$. For large j we have that $\cos^{2j}(\phi/2) \sim \exp[-(j/4)\phi^2]$, while $\sin^{2j}(\phi/2) \sim \exp[-(j/4)(\phi - \pi)^2]$. Therefore, the phase probability distribution is a Gaussian with width scaling as $\sigma = \sqrt{2}/\sqrt{N_T}$ (see Fig.(4,B)). The NOON state with a Mach-Zehnder interferometer, assuming a complete “a priori” ignorance of the real phase shift, achieves a shot noise sensitivity. A previous analysis [30], di-

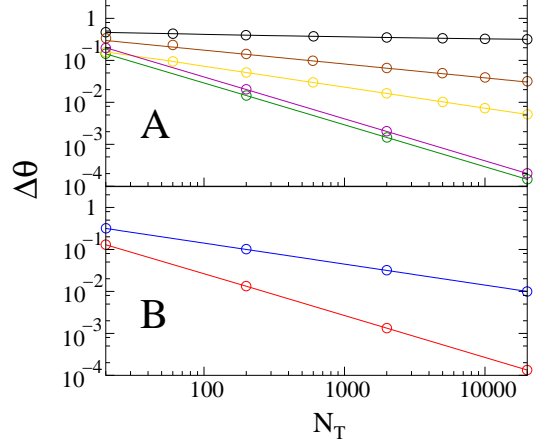


FIG. 4: Phase uncertainty $\Delta\theta$ as a function of N_T . A) Twin-Fock state: (from the top line to the bottom one) black line, σ for $p = 1$ ($\sigma \approx 1/\ln N_T$); brown line, 68% confidence for $p = 1$ ($\Delta\theta = 0.8/N_T^{1/3}$); yellow line, σ for $p = 2$ ($\sigma \approx 1/N_T^{1/2}$); violet line, σ for $p = 4$ ($\sigma = 4/N_T$); green line: 68% confidence for $p = 2$ ($\Delta\theta = 2.91/N_T$). B) Phase sensitivity of the NOON state (blue line), $\sigma = \sqrt{2}/N_T$, independent of p , and of the twin $m = 1$ state (red line), $\Delta\theta = 2.67/N_T$ at $p = 1$. Points are numerical data.

rectly addressing current experiments [16, 17], showed that NOON states and balanced homodyne detection can provide phase sensitivity scaling as $\sim 1/N_T^{3/4}$.

Mach-Zehnder interferometer with a twin $m = 1$ state. We have previously shown how the large tails of the twin-Fock probability prevent the sensitivity to scale as $\sim 1/N_T$. To kill these tails, and enhance the central peak of the distribution, we have to statistically analyze several independent measurements. Here we follow a different strategy: is it possible to cancel out the tails by destructive interference? The answer is positive. The tails of the probability distribution $P(\phi|j, 0) \sim |d_{m,m}^j(\phi) + d_{m,-m}^j(\phi)|^2$ obtained with the input state (1) cancel out when m is odd. This special feature, not present in the twin-Fock state, is a consequence of the interference between the two matrices in $P(\phi|j, 0)$, and it can be easily recognized by expanding the Jacobi Polynomials [31]. Since the cancellation has a precision of the order of $O[(j - m)^{-3}]$ we can conclude that, among this class of states, the case $m = 1$ offers the highest performance. States with larger odd values of m do not benefit of an equally efficient destructive interference. With semi-integer values of m , as for the Yurke state ($m = 1/2$) [4], there is no cancellation at all, so we expect that a $1/N_T$ scaling can be reached only with $p > 1$ as in the twin-Fock case. In Fig.(2,C) we plot the amplitude $d_{1,1}^j(\phi)$ (red line), and $d_{1,-1}^j(\phi)$ (blue line), for $N = 40$ particles ($j = 20$). Outside the central region the two functions oscillate out of phase and, when summed to calculate $P(\phi|j, 0)$, interfere destructively. The phase probability distribution is shown in Fig.(2,D). The comparison with the Fock state distribution (Fig. (2,A)), for the same total number of particles, highlights the strong

reduction of the tails. As a consequence, for a given confidence, we reach the minimum with a smaller number of measurements with respect to the Fock state case, see Fig. (3,B). The sensitivity of the twin- m states is:

$$\Delta\theta_{68\%} = \frac{2.67}{N_T}, \quad p = 1; \quad \Delta\theta_{95\%} = \frac{5.376}{N_T}, \quad p = 1. \quad (5)$$

The scaling of $\Delta\theta_{68\%}$ as a function of N_T is shown in Fig.(4,B). The main advantage is, apart from the smaller prefactors, that the $1/N_T$ scaling is reached with a single ($p = 1$) measurement for both the 68% and the 95% confidence, as compared with the $p = 2$ and $p = 3$ measurements requested with Fock states, respectively. For higher confidences, the required number of measurements is shown in Fig.(3,B) [32, 33].

Conclusions. There are not known protocols for the direct measurement of phase shifts. Phases can only be inferred after the measurement of a different (phase dependent) observable. What is the highest achievable sensitivity? Within a rigorous Bayesian analysis we find that $|\psi\rangle_{inp} = 1/\sqrt{2}(|N_T/2+1\rangle|N_T/2-1\rangle + |N_T/2-1\rangle|N_T/2+1\rangle)$, gives the *highest sensitivity* $C_{68\%} = 2.67/N_T$, with a *single* interferometric measurement. Different input quantum states can achieve the Heisenberg scaling $\sim 1/N_T$ but with higher prefactors and at the price of a statistical analysis of *two or more* independent Mach-Zehnder measurements.

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 - [32] The possibility to estimate a phase with a single measurement has some important advantages. Indeed, when the external perturbation has a short life-time, the need to perform several independent measurements might require the building of an equal number of interferometers.
 - [33] We have extended our analysis to consider classes of states different from Eq.(1). We have always found that the twin $m = 1$ state provides the best performance, and, on the ground of some general arguments, we believe that this remains true with any possible two-mode input quantum state.